Physical model for the L and H regimes in a tokamak

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The structure of the particle and energy fluxes in the plasma is analyzed in an effort to explain the L and H regimes in a tokamak. A pinch in the particle flux is turned on or off, depending on the extent to which the plasma is collisional. The result is a transition from the L regime to the H regime. The $OH \to L$ transition occurs because a new turbulent transport coefficient comes into play.

Although it has been several years since Wagner *et al.*[1] found experimentally that there are two plasma confinement regimes (the L and H regimes) in a tokamak, we do not yet have a unified model for the physics of the transition to the L and H regimes. In the L regime the confinement is degraded in comparison with the ohmic-heating regime. The transition to the H regime correlates best with a cutoff of the particle flux at the plasma periphery and the presence of a divertor configuration of the magnetic field.[2] In the present letter we analyze a possible form of the particle and energy fluxes and also expressions for the transport coefficients which follow from the theory. We begin with a look at particle transport, which is controlled by the electrons, which have the smaller diffusion coefficient. The ion transport adjusts to accommodate the electron transport by virtue of the ambipolar nature of the process. It can be shown[3,4] that in the collisionless regime the particle flux Γ which is caused by the interaction of the untrapped and trapped electrons with plasma waves is

$$\Gamma=\Gamma_u+\Gamma_{tr}=-D_u\left(\frac{\partial n}{\partial r}-a_u\frac{n}{T}\frac{\partial T}{\partial r}\right)-\epsilon^{1/2}D_{tr}\left(\frac{\partial n}{\partial r}-a_{tr}\frac{n}{T}\frac{\partial T}{\partial r}\right). \qquad (1)$$

Here u and tr specify the untrapped and trapped particles, and the notation is otherwise standard. The numerical coefficients are $a_u \approx 1/2$ and $a_{tr} \approx 1$. The pinch velocity of the untrapped particles, $V_u = (D_u/2T)\partial T/\partial r$, appears here because of the one-dimensional nature of the distribution function of the untrapped electrons and the assumption that the phase velocity of the waves in low $(\omega/k_{\parallel} \approx V_{\parallel} \ll V_{Te})$, when we have $f_e \approx n/T^{1/2}$. In the case of the trapped particles the pinch velocity $V_{tr} = D_{tr}a_{tr}(1/T)(\partial T/\partial r)$ also arises from a differentiation of the distribution function f_e, which is now two-dimensional. Furthermore, a_{tr} depends on the wave spectrum in the plasma,[4] varying over the range 0.5–1.5.

While the fluxes of the untrapped and trapped particles, Γ_u and Γ_{tr}, respectively, are superficially similar, the diffusion coefficients involved in them, D_u and D_{tr}, have completely different structures. For D_u the quasilinear theory predicts[3]

$$D_u = C^2 V_{Te}\epsilon^{\alpha}(\omega_{pe}{}^2/qR), \qquad (2)$$

in agreement with the Merezhkin-Mukhovatov scaling[5] with $\alpha = 1.75$. For trapped particles, an expression for D_{tr} was derived in the quasilinear theory in Ref. 4.

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The diffusion coefficients D_u and D_{tr} and the structure of flux (1) were derived in a collisionless regime in the quasilinear approximation. As the wave level in the plasma increases, a highly nonlinear turbulent regime sets in, and the structure of the particle flux changes: The pinch disappears, and there is a change in the diffusion coefficient, which can be estimated from

$$D_d = C_s \hat{\rho}_s^{\,2} q(1+q^2/2)^{0,5} / (r_n S \epsilon_n). \tag{3}$$

Here $r_n = n/(\partial n/\partial r)$, $\epsilon_n = r_n/R$, $C_s = (T_e/M)^{1/2}$, $\rho_s = C_s/\omega_{Bi}$, $S(r/q)\partial q/\partial r$ is the shear, and q is the safety factor in the tokamak. Estimate (3) follows from the expression $D = \Sigma_k \gamma_k |\xi_\chi|_k^2 \sim \gamma/k_\chi^2$, where ξ_χ is the radial displacement of an element of the plasma. This estimate is complicated by the circumstance that the convection cells in a plasma with a drift turbulence take the form of not only single-scale "round" vortices with a size $\sim\rho_s$ but also two-scale formations which are narrow ($\sim\rho_s$) and highly elongated ("rivers"), in which the plasma is transported over a distance L which is far larger than ρ_s during the decorrelation time $\sim\gamma^{-1}$. Working from the results of the numerical calculation of Ref. 8, we would naturally as k_χ^{-1} the largest displacement L, which is associated with the length of a river: $k_\chi^2 \approx (\gamma/\omega_*^2)|k_\parallel' C_s|/\rho_s$. We then find estimate (3) (ω_* is the drift frequency). Incorporating the toroidal curvature leads to an additional intensification of the convection, by a factor of $(1+q^2/2)^{0.5}$ The estimate of the coefficient D_d is described in more detail in Ref. 6. The coefficient D_d has the important property that at $q>1$ it is no longer dependent on the strength of the toroidal magnetic field B_T. Instead, a dependence on the field of the current, B_ϑ arises, i.e., a dependence on the current I. The scaling corresponding to D_d, $\tau_E \propto f(n) I^{0.8} P^{-0.6}$, where P is the power deposition, agrees with the experimental scaling found for the confinement in the L regime,[7] except for the independence. Dimensional estimate without a pinch and without the profiles yield $f(n) \propto n^{0.6}$. A numerical simulation reveals $f(n) \approx$ const, in agreement with the experiments of Ref. 7. This discrepancy can be explained quite easily on the basis that a pinch in the particle flux Γ is turned on and off (as we will discuss below); this is actually a divergence of the local scaling and the global scaling.

We thus find the following structure for particle flux Γ:

$$\Gamma = -(D_u + D_d)\partial n/\partial r + a_1 D_u (n/T)\partial T/\partial r, \quad a_1 \approx 1. \tag{4}$$

With decreasing mean free path λ, the constant a_1 becomes a function of n, T, and r; if λ is sufficiently small, the pinch disappears. The physical reason for the turning off of the pinch is that as the collision rate ν_{eff} increases, the waves begin to interact not exclusively with some distinct group of particles in phase space but with essentially all of the particles in a given volume element of the plasma, because of a broadening in the reasonances. The structure of the phase space is no longer important in this case. We simulated a collisional turning off a pinch by means of

$$a_1 \neq 0, \text{ for } \lambda > \lambda_*, \quad a_1 = 0, \text{ for } \lambda < \lambda_*, \tag{5}$$

where $\lambda = V_{Te}/\nu_{\text{eff}}$ is the mean free path of the particles, $\lambda_* = qRF(\epsilon)Z_{\text{eff}}^2$, and $F(\epsilon)$ incorporates the contribution from the untrapped ($F_u \approx 1$) and trapped ($F_{tr} \approx \epsilon^{1.5}$)

particles. Condition (5) plays a key role in the model for the transition to the H regime.

By analogy with (1), the energy flux generally contains two terms, which are proportional to gradients $\partial T/\partial r$ and $\partial n/\partial r$:

$$Q = -1.5\chi(n\partial T/\partial r - a_2 T\partial n/\partial r). \qquad (6)$$

Here $|a_2| \approx 1$. Writing an expression for $\partial n/\partial r$ from (4), we find

$$Q = -\frac{3}{2}\chi(1 - a_1 a_2 D_u/D)n\frac{\partial T}{\partial r} + (-\frac{3}{2}a_2\frac{\chi}{D})T\Gamma = -\frac{3}{2}\chi_{eff} n\frac{\partial T}{\partial r} + \frac{5}{2}T\Gamma, \qquad (7)$$

where $D = D_u + D_d$. The first term in (7) describes the ordinary heat flux, while the second is associated with a transport of energy directly by the particles. It follows that we have $a_2 < 0$; there is no net pinch in the energy flux. The coefficient of $T\Gamma$ in (7) is

$$1.5(\chi_{eff}/D)|a_2| / (1 + a_1|a_2| D_u/D) \approx 1.5 - 3,$$

and it is usually assumed to be 5/2. This assumption seems arbitrary in view of the uncertainty regarding the values of the ratio $\chi_{\rm eff}/D$ and also a_1 and a_2.

In accordance with transport model (4), (7), we can outline the following scenario for the discharge: In the ohmic-heating regime the plasma temperature is not very high, and the transport coefficient D_u and the corresponding $\chi_{\rm eff} \sim D_u$ are predominant, since the relation $D_d \ll D_u$ holds. The pinch is turned on first in the central part of the plasma column and is not present at the periphery of the plasma. In this regime the Merezhkin-Mukhovatov scaling holds,[5] complicated somewhat by the presence of a pinch. During auxiliary heating the plasma temperature rises, and the coefficient $D_d \propto T^{3/2}$ becomes greater than $D_u \propto T^{0.5}/n$. This degradation of the confinement corresponds to a transition to the L regime, in which, as we mentioned above, the scaling of the confinement is controlled by the coefficient D_d. In the L regime the pinch in the particle flux exists over nearly the entire cross section of the column, except at the periphery, where the plasma is still rather cool and contaminated ($Z_{\rm eff} > 1$), so the second inequality in (5) holds. If, on the other hand, auxiliary heating is carried out at the periphery, and/or the periphery is purified by means of a divertor (if $Z_{\rm eff}$ is lowered), the pinch will also turn on at the periphery, according to (5); it will block the particle flux Γ and substantially improve the confinement (this is the transition to the H regime). This scenario of the transitions to the L and H regime is confirmed by the results of our numerical simulation based on the particle and energy balance equations, (4) and (7), with (2), (3), and (5).

In summary, we link the existence of two plasma confinement regimes (the L and H regimes) in a tokamak with manifestations of two major effects: (1) the turning on and turning off of a pinch in the particle flux Γ, depending on the extent to which the plasma is collisional, especially at the periphery of the plasma column, and (2) a change in the relative roles played by the collisionless quasilinear loss mechanism ($D_u \propto T^{0.5}$) and the turbulent loss mechanism ($D_d \propto T^{1.5}$), depending on the plasma temperature.

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